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Spontaneous R-Parity violation bounds

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Abstract

We investigate bounds from tree-level and one-loop processes in generic supersymmetric models with spontaneous R-parity breaking in the superpotential. We analyse the bounds from a general point of view. The bounds are applicable both for all models with spontaneous R-parity violation and for explicit bilinear R-parity violation based on general lepton-chargino and neutrino-neutralino mixings. We find constraints from semileptonic B, D and K decays, leptonic decays of the μ and τ, electric dipole moments, as well as bounds for the anomalous magnetic moment of the muon.

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1 Introduction

While supersymmetry appears to be the best scenario for physics beyond the Standard Model (SM), most of the early studies have been made in the context of the minimal supersymmetric standard model (MSSM), the supersymmetric analog of the Standard Model. This model assumes conservation of a discrete symmetry called the R-parity, which is related to baryon number, lepton number and spin through $R = (-1)^{(3B+L+2S)}$. Under this symmetry all the Standard Model particles are R-even, while their superpartners are R-odd. With this assumption the supersymmetric particles must be pair produced, every supersymmetric particle decays into another and the lightest of them is stable. Although the MSSM has some attractive features, neither gauge invariance nor supersymmetry require R-conservation, so within the MSSM R-conservation is imposed as a symmetry of the model. The most general supersymmetric extension of the Standard Model containing explicit R-violating interactions has received a lot of attention lately as one of the possible explanations for neutrino masses and oscillations. Numerous detailed analyses of the explicit (bilinear or trilinear) Yukawa couplings have appeared in the literature along with constraints on these couplings [1, 2].

Less attention has been given to the possibility that R is an exact symmetry of the Lagrangian, but broken spontaneously through the Higgs mechanism. This would occur through scalar neutrinos acquiring non-zero vacuum expectation values:

$$\langle \tilde{\nu}_{L_i} \rangle \neq 0 \quad ; \quad \langle \tilde{\nu}_{R_i} \rangle \neq 0 \tag{1}$$

Such a breaking is natural in scenarios beyond the MSSM, such as for instance in the left-right supersymmetric model (LRSUSY) where spontaneous R-parity breaking avoids the existence of a charge-violating minimum in the superpotential [3]. If spontaneous R-parity violation occurs in the absence of any additional gauge symmetries, it will lead to the existence of a physical massless Nambu-Goldstone boson, called a Majoron [4]. Phenomenological studies of spontaneous R-parity breaking have mostly concentrated on the experimental consequences of the Majoron, the most significant of which is the increase of the invisible Z^0 width by an amount equivalent to 1/2 of a light neutrino family [5].

Another phenomenologically interesting consequence of the spontaneous R-symmetry breaking is that it introduces new terms in the Lagrangian, as compared to the Yukawa couplings of the explicit trilinear R-parity breaking. In particular it introduces interactions

with gauginos which have the feature that they are sfermion-mass independent. Precision measurements of rare processes put rather strong constraints on such spontaneous R-breaking terms. These interactions appear in explicit bilinear breaking as soft dimensionful Higgs-lepton superfield mixing parameter.

In this work we study the phenomenological implications of spontaneous R-parity breaking in the supersymmetric Lagrangian. We analyse the general form of the gaugino-higgsino-lepton mixing and set the most general bounds on the mixing elements based on rare tree-level and one-loop level processes. The advantage of setting bounds on the mixing elements lies in their generality: they apply to any supersymmetric model with spontaneous R-parity breaking, or even to a SUSYGUT scenario with an enriched gauge sector. Our paper is organized as follows: we describe and parametrize spontaneous R-parity breaking in section 2, then discuss tree-level constraints in section 3, one-loop level constraints in section 4, before reaching our conclusion in section 5.

2 Spontaneous R-parity breaking

Specific superpotentials can be designed to violate R-parity and lepton number spontaneously [6]. We concentrate here on the phenomenological consequences. As a consequence of spontaneous R-parity breaking, the sneutrino fields $\tilde{\nu}_i$ acquire non-zero vacuum expectation values $\langle\tilde{\nu}_i\rangle \neq 0$. In order to have the spontaneous breaking of the R-parity, new fields have to be added to the MSSM spectrum. In order to set spontanous R parity breaking in perspective, we outline briefly the main features of two models present in the literature.

In the model proposed in [7], a superpotential which conserves total lepton number and R-parity is constructed. Additional fields (Φ, ν_i^c, S_i) are introduced, which are singlets under $SU(2)_L \times U(1)_Y$ and which carry a conserved lepton number assigned as $(0, -1, 1)$. The superpotential has the form:

$$W = h_u Q H_u U + h - d Q H_d D + (h_0 H_u H_d - \epsilon^2)\Phi + h_\nu L H_u \nu^c + h \Phi S \nu^c. \tag{2}$$

R-parity is broken spontaneously, if one or more of the singlets have a vev: $v_R = \langle\tilde{\nu}_\tau^c\rangle$, $v_S = \langle S_\tau\rangle$ and $v_L = \langle\tilde{\nu}_{L\tau}\rangle$. The vev of the isodoublet Higgs will drive the electroweak symmetry breaking and allow fermion masses in the usual fashion. The bounds on the sneutrino vev's have been considered in the bilinear R-parity breaking model. From the

Superkamiokande data, the constraints for the vev's are obtained as [8]

$$
\begin{aligned}
\langle\tilde{\nu}_e\rangle/\sqrt{M_{SUSY}/100\,\mathrm{GeV}} &\leq 90\,\mathrm{keV}, \\
76\,\mathrm{keV} \leq \quad \langle\tilde{\nu}_{\mu,\tau}\rangle/\sqrt{M_{SUSY}/100\,\mathrm{GeV}} &\leq 276\,\mathrm{keV}.
\end{aligned}
\tag{3}
$$

In [6] it has been found that $t-b-\tau$ unification is allowed for $\langle\tilde{\nu}_\tau\rangle \lesssim 5$ GeV and $b-\tau$ unification is allowed for $\langle\tilde{\nu}_\tau\rangle \lesssim 50$ GeV.

In the left-right supersymmetric model, based on the gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, the R-parity is a discrete subgroup of the $U(1)_{B-L}$. It is possible that in the process of spontaneous symmetry breaking this model develops a minimum which violates R-parity; indeed in some versions of LRSUSY this breaking is inevitable [3, 9]. In the minimal version of this model, the most general gauge invariant superpotential must contain, in addition to the usual left and right-handed quark and lepton doublets, two Higgs bidoublets Φ_u and Φ_d and four Higgs triplet superfields $\Delta_{L,R}$ and $\delta_{L,R}$. The superpotential corresponding to this minimal field content is:

$$
\begin{aligned}
W_{min} &= Q^T i\tau_2(h_{\Phi_u}\Phi_u + h_{\Phi_d}\Phi_d)Q^c + L^T i\tau_2(h_{\Phi_u}\Phi_u + h_{\Phi_d}\Phi_d)L^c \\
&+ h_\Delta(L^T i\tau_2\delta_L L + L^{cT} i\tau_2\Delta_R L^c) + \mu_{ij}Tr(i\tau_2\Phi_i^T i\tau_2\Phi_j) + \mu_\Delta(\Delta_L\delta_L + \Delta_R\delta_R)
\end{aligned}
\tag{4}
$$

In the minimal model, breaking parity spontaneously at the renormalizable level is always accompanied by spontaneous R parity breaking. This may be cured by adding more fields to the theory [3]. One is left with a relatively low $SU(2)_R$ breaking scale, with the spontaneously broken R-parity driven by $\sigma_R = \langle\tilde{\nu}^c\rangle \neq 0$.

In this work we will not assume any particular model for the breaking. Instead we will study interactions typical for this class of models. In what follows, we will present the formulas with the MSSM particle content. This would effectively be the case if the other fields in the model decouple. However, the formalism described can be extended straightforwardly to richer matter/gauge sectors, such as e.g, MSSM with right-handed neutrinos, where both the left-handed and the right-handed sneutrinos can acquire a vev, or left-right model, where the number of gauginos and higgsinos is larger than in the MSSM.

Within the minimal field content, the chargino-lepton mixing matrix becomes 5×5 matrix and the neutralino-neutrino matrix a 7×7 matrix. The mass eigenstate fields can be written as:

$$
\Psi_i^0 = N_{ij}\Psi_j'^0, \qquad \Psi_i^+ = V_{ij}\Psi_j'^+, \qquad \Psi_i^- = U_{ij}\Psi_j'^- \tag{5}
$$

for, respectively, the neutral and charged fields, where the weak eigenstates are:

$$\Psi_j^{\prime 0T} = (-i\lambda', -i\lambda_3, \tilde{H}_1^0, \tilde{H}_2^0, \nu_i), \quad i = e, \mu, \tau \tag{6}$$

$$\Psi_j^{\prime -T} = (-i\lambda_-, \tilde{H}_1^-, e_L^-, \mu_L^-, \tau_L^-), \tag{7}$$

$$\Psi_j^{\prime +T} = (-i\lambda_+, \tilde{H}_2^+, e_L^+, \mu_L^+, \tau_L^+). \tag{8}$$

are the $U(1)_Y$ gauginos, $\lambda_{3,+,-}$ are the $SU(2)_L$ gauginos, and $\tilde{H}$ are the higgsinos. To extend to more complicated particle contents one needs to add more gauginos and higgsinos in Eqs. (6)-(8).

The relevant part of the interaction Lagrangian becomes, for quarks-squarks-charginos:

$$\begin{aligned}
\mathcal{L}_{q\tilde{q}'\Psi^+} &= -g \sum_i \Big\{ \bar{\Psi}_i^+ [(U_{i1}^* P_L - \frac{m_{u_k} V_{i2}}{\sqrt{2} M_W \sin\beta} P_R) u_k \tilde{d}_{Lk}^* \\
&- \frac{m_{d_k} U_{i2}^*}{\sqrt{2} M_W \cos\beta} P_R u_k \tilde{d}_{Rk}^*] + \bar{u}_k (U_{i1} P_R - \frac{m_{u_k} V_{i2}^*}{\sqrt{2} M_W \sin\beta} P_L) \Psi_i^+ \tilde{d}_{Lk} \\
&- \frac{m_{d_k} U_{i2}}{\sqrt{2} M_W \cos\beta} P_L \Psi_i^+ \tilde{d}_{Rk}] + \bar{\Psi}_i^{+c} (V_{i1}^* P_L - \frac{m_{d_k} U_{i2}}{\sqrt{2} M_W \cos\beta} P_R) d_k \tilde{u}_{Lk}^* \\
&- \frac{m_{u_k} V_{i2}^*}{\sqrt{2} M_W \sin\beta} P_R d_k \tilde{u}_{Rk}^*] + \bar{d}_k (V_{i1} P_R - \frac{m_{d_k} U_{i2}^*}{\sqrt{2} M_W \cos\beta} P_L) \Psi_i^{+c} \tilde{u}_{Lk} \\
&- \frac{m_{u_k} V_{i2}}{\sqrt{2} M_W \sin\beta} P_L \Psi_i^{+c} \tilde{u}_{Rk}] \Big\}
\end{aligned} \tag{9}$$

and for quarks-squark-neutralinos:

$$\begin{aligned}
\mathcal{L}_{q\tilde{q}\Psi^0} &= -\sqrt{2} \sum_j \Big\{ \bar{u}_k \{ [e e_u N_{j1} + \frac{g}{\cos\theta_W} (1/2 - e_u \sin^2\theta_W N_{j2})] P_R \Psi_j^0 \tilde{u}_{Lk} \\
&+ \frac{g m_{u_k}}{2 M_W \sin\beta} N_{j4}^* P_L \Psi_j^0 \tilde{u}_{Lk} - [e e_u N_{j1}^* - (\frac{g e_u \sin^2\theta_W}{\cos\theta_W}) N_{j2}^*)] P_L \Psi_j^0 \tilde{u}_{Rk} \\
&+ \frac{g m_{u_k}}{2 M_W \sin\beta} N_{j4} P_R \Psi_j^0 \tilde{u}_{Rk} \} \\
&+ \bar{d}_k \{ [e e_d N_{j1} - \frac{g}{\cos\theta_W} (1/2 + e_d \sin^2\theta_W N_{j2})] P_R \Psi_j^0 \tilde{d}_{Lk} \\
&+ \frac{g m_{d_k}}{2 M_W \cos\beta} N_{j3}^* P_L \Psi_j^0 \tilde{d}_{Lk} - [e e_d N_{j1}^* - (\frac{g e_u \sin^2\theta_W}{\cos\theta_W}) N_{j2}^*)] P_L \Psi_j^0 \tilde{d}_{Rk} \\
&+ \frac{g m_{d_k}}{2 M_W \cos\beta} N_{j3} P_R \Psi_j^0 \tilde{d}_{Rk} \} \Big\}.
\end{aligned} \tag{10}$$

Similar expressions are obtained for the lepton-slepton interactions. Note that in both parts of the Lagrangian, Eqs. (9) and (10), we get interactions which do not depend on mass, but depend on the mixing with the gaugino. This is a major difference compared to the explicit

trilinear R-parity breaking, in which only those terms which depend on mass and also on the mixing with higgsino are present.

Assuming that the new fields, which transform the potential so that the R-parity breaks, do not mix with the MSSM higgsinos and gauginos, the mass matrices for the spontaneous R-parity breaking and explicit bilinear R-parity breaking are very similar. With this assumption, the mass matrices become ($i, j = e, \mu, \tau$):

$$(M)_{\Psi^\pm} \;=\; \left(\begin{array}{ccc} M_2 & \frac{1}{\sqrt{2}} g v_u & 0 \\ \frac{1}{\sqrt{2}} g v_d & \mu & \frac{1}{\sqrt{2}} h_i \langle \tilde{\nu}_{Li} \rangle \\ \frac{1}{\sqrt{2}} g \langle \tilde{\nu}_{Lj} \rangle & h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle & \frac{1}{\sqrt{2}} h_i v_d \delta_{ij}, \end{array} \right) \tag{11}$$

for the chargino-lepton (where h_i are Yukawa couplings from the R-conserving Lagrangian $h_i L_i H_1 E_i$), and:

$$(M)_{\Psi^0} \;=\; \left(\begin{array}{ccccc} M_1 & 0 & \frac{1}{\sqrt{2}} g' v_d & \frac{1}{\sqrt{2}} g' v_u & \frac{1}{\sqrt{2}} g' \langle \tilde{\nu}_{Li} \rangle \\ 0 & M_2 & \frac{1}{\sqrt{2}} g v_d & \frac{1}{\sqrt{2}} g v_u & \frac{1}{\sqrt{2}} g' \langle \tilde{\nu}_{Li} \rangle \\ \frac{1}{\sqrt{2}} g' v_d & \frac{1}{\sqrt{2}} g v_d & 0 & -\mu & 0 \\ \frac{1}{\sqrt{2}} g' v_u & \frac{1}{\sqrt{2}} g v_u & -\mu & 0 & -h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle \\ \frac{1}{\sqrt{2}} g' \langle \tilde{\nu}_{Li} \rangle & \frac{1}{\sqrt{2}} g \langle \tilde{\nu}_{Li} \rangle & 0 & -h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle & 0 \end{array} \right) \tag{12}$$

for the neutralino-neutrino. The lightest mass eigenstates obtained by diagonalizing the mass matrices correspond to the neutrinos and charged leptons. By rotating the MSSM Lagrangian to the new mass eigenstates one obtains new lepton-flavor violating interactions. The mixing matrices induced by the $\mathcal{L}_{LH}$ are listed below. For the charged sector the matrices are [10]:

$$(U)^* \;=\; \left(\begin{array}{cc} U_R(1 - \frac{1}{2}\xi^{L^T}\xi^{L*}) & -V_L \xi^{L*} \\ U_R \xi^{L^T} & V_L(1 - \frac{1}{2}\xi^{L*}\xi^{L^T}) \end{array} \right), \tag{13}$$

$$(V)^\dagger \;=\; \left(\begin{array}{cc} (1 - \frac{1}{2}\xi^{R^T}\xi^{R*}) U_L^\dagger & \xi^{R*} U_L^\dagger \\ -\xi^{R^T} V_R^\dagger & (1 - \frac{1}{2}\xi^{R*}\xi^{R^T}) V_R^\dagger \end{array} \right). \tag{14}$$

In the neutral fermion sector the mixing matrix is:

$$(N)^* \;=\; \left(\begin{array}{cc} N^{0*}(1 - \frac{1}{2}\xi^\dagger \xi) & -V^{(\nu)^T}\xi \\ N^{0*}\xi^\dagger & V^{(\nu)^T}(1 - \frac{1}{2}\xi\xi^\dagger), \end{array} \right) \tag{15}$$

In the above equations the parameters $\xi^L,\ \xi^R,\ \xi$ represent mixing beween the MSSM sector matrices (corresponding to the matrices $U_L,\ U_R$ for chargino and N^0 for the neutralino in MSSM) and the lepton ($V_L,\ V_R$) or neutrinos ($V^{(\nu)}$) mixing matrices. The relationship

between these matrices and the MSSM matrices is:

$$U_R M_{\chi^\pm} U_L^\dagger = \mathrm{Diag}\{M_{\chi_i^\pm}\}, \tag{16}$$

$$V_L M^{(l)} V_R^* = \mathrm{Diag}\{m_{l_i}\}, \tag{17}$$

$$N^{0*} M_{\chi^0} N^{o\dagger} = \mathrm{Diag}\{M_{\chi_i^0}\}, \tag{18}$$

$$V^{(\nu)^T} m_{eff} V^\nu = \mathrm{Diag}\{m_{\nu_e}, m_{\nu_\mu}, m_{\nu_\tau}\}, \tag{19}$$

where the mixing parameters $\xi^L,\ \xi^R,\ \xi$ are, for $i = 1, 2, 3$:

$$\xi_{i1}^{*L} = \frac{g_2}{\sqrt{2}\det M_{\chi^\pm}} \Lambda_i, \tag{20}$$

$$\xi_{i2}^{*L} = \frac{h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle}{\mu} - \frac{g_2 \sin\beta M_W}{\mu \det M_{\chi^\pm}} \Lambda_i, \tag{21}$$

$$\xi^{*R} = M^{(l)\dagger} \xi^{*L} (M_{\chi^\pm}^{-1})^T, \tag{22}$$

$$\xi_{i1} = \frac{g_1 M_2 \mu}{2 \det M_{\chi^0}} \Lambda_i, \tag{23}$$

$$\xi_{i2} = -\frac{g_2 M_1 \mu}{2 \det M_{\chi^0}} \Lambda_i, \tag{24}$$

$$\xi_{i3} = \frac{h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle}{\mu} + \frac{g_2 (M_1 + \tan^2\theta_W M_2) \sin\beta M_W}{2 \det M_{\chi^0}} \Lambda_i, \tag{25}$$

$$\xi_{i4} = -\frac{g_2 (M_1 + \tan^2\theta_W M_2) \cos\beta M_W}{2 \det M_{\chi^0}} \Lambda_i. \tag{26}$$

Here $\chi^\pm$ and χ^0 denote the MSSM charginos and neutralinos. In the above expressions $\Lambda_i = \mu \langle \tilde{\nu}_{Li} \rangle - \langle H_1 \rangle h_{\nu_{ij}} \langle \tilde{\nu}_{Rj} \rangle$ is a measure of the misalignment and is small, but must be essentially nonzero for neutrinos to have a mass. Also, in the above μ is the MSSM bilinear Higgs coupling, and $\det M_{\chi^\pm} = M_2 \mu - \sin 2\beta M_W^2$ is the determinant of the MSSM chargino mass matrix; $\det M_{\chi^0} = M_W^2 \mu \sin 2\beta (M_1 + M_2 \tan^2\theta_W) - M_1 M_2 \mu^2$ is the determinant of the MSSM neutralino mass matrix.

The spontaneous R-parity breaking is driven by $\langle \tilde{\nu}_{Li} \rangle$ and $\langle \tilde{\nu}_{Ri} \rangle$. The mass matrices in bilinear R-parity breaking can be obtained by making the substitution $h_{\nu ij} v_{Rj} \to \epsilon_j$, which makes it easy to read the mixing matrices from one case to the other.

The parametrization we presented above is only one of the ones available in the literature. Others exists, most notable the single vev parametrization [11]. In the next sections we will present bounds on the matrix elements themselves coming from phenomenological constraints, which are independent of any parametrization chosen.

3 Tree-level bounds on matrix elements from rare decays

3.1 Semileptonic B, D and K decays

In this section we investigate constraints arising from rare decays of the B, D and K mesons, as well three-body lepton number violating decays of the μ and τ, all of which can occur at tree-level and all of which put bounds of spontaneous R-parity violating matrix mixing elements. Since all these decays occur at tree-level through an exchange of a scalar fermion, we will employ throughout the notation:

$$q_i = \left(\frac{100\ GeV}{m^2_{\tilde{q}_i}}\right)^2,\ l_i = \left(\frac{100\ GeV}{m^2_{\tilde{l}_i}}\right)^2,\ n_i = \left(\frac{100\ GeV}{m^2_{\tilde{\nu}_i}}\right)^2 \tag{27}$$

with $q = u, d$, and $i = 1, 2, 3$ represents the three quark families, $l_i = e, \mu, \tau$ and $\nu_i = \nu_e, \nu_\mu, \nu_\tau$.

We first analyse the semileptonic decays of the K, D and B-mesons. The effective Lagrangians relevant for semileptonic decays of the B-mesons are:

$$\mathcal{L}_{eff}(b \to q l_i \nu_j) = -K_{qb}\frac{4G_F}{\sqrt{2}}[\mathcal{A}^q_{ij}(\bar{q}\gamma^\mu P_L b)(\bar{l}_i\gamma_\mu P_L \nu_j) - \mathcal{B}^q_{ij}(\bar{q}P_R b)(\bar{l}_i P_L \nu_j)], \tag{28}$$

and

$$L_{eff}(b \to q l^+_i l^-_j) = -K_{qb}\frac{4G_F}{\sqrt{2}}[\mathcal{C}^q_{ij}(\bar{q}\gamma^\mu P_L b)(\bar{l}_i\gamma_\mu P_L l_j) - \mathcal{D}^q_{ij}(\bar{q}P_R b)(\bar{l}_i P_L l_j)], \tag{29}$$

where K is the CKM matrix. The Lagrangian is similar for all the semileptonic decays with the appropriate substitutions for the b quark. The leptonic branching ratios for the processes $b \to e\nu X$ and $b \to \mu\nu X$ measured by the L3 Collaboration [12] are:

$$\begin{aligned} BR(b \to e\nu X) &= (10.89 \pm 0.55) \times 10^{-2}, \\ BR(b \to \mu\nu X) &= (10.82 \pm 0.61) \times 10^{-2}. \end{aligned} \tag{30}$$

These decay processes can occur at tree-level through either $\tilde{d}$ and $\tilde{b}$ or $\tilde{u}$ and $\tilde{t}$ exchanges. They set bounds on both the higgsino and the gaugino couplings in spontaneous R-parity violating models.

The branching ratios for semileptonic decays into charged leptons will set bounds on the chargino-lepton mixing elements only. The present measurements of the branching ratios of

the $b \to s l_j^+ l_i^-$ give the following upper bounds (at 90% C.L.) [13]

$$\begin{aligned} BR(b \to s e^+ e^-) &< 5.7 \times 10^{-5}, \\ BR(b \to s \mu^+ \mu^-) &< 5.8 \times 10^{-5}, \\ BR(b \to s e^\pm \mu^\mp) &< 2.2 \times 10^{-5}. \end{aligned} \tag{31}$$

The experimental bounds on the first two are almost one order of magnitude larger than the SM expectation: the last process is forbidden in SM because of separate conservation of each lepton flavor number. The bounds obtained are listed in Table 1. They involve products of neutralino-neutrino and chargino-lepton mixing matrices.

Table 1: Analytic bounds on mixing matrices of chargino-leptons $U_{ij},\ V_{ij}$ (with $j = 3,4,5$ corresponding to e, μ and τ) and neutralino-neutrinos N_{ij} (with $j = 5,6,7$ corresponding to ν_e, ν_μ and ν_τ) from B rare semileptonic decays.

gaugino type	higgsino type	bound	process
$-\sqrt{2}\|K_{td}\|\,\{e[V_{31}e_dN^*_{i1}(d_1+d_3)$ $-V^*_{31}e_uN_{i1}(u_1+u_3)]$ $-g\left[\frac{(1/2+e_d\sin^2\theta_W)(d_1+d_3)}{\cos\theta_W}V_{31}N^*_{i2}\right.$ $\left.-\frac{(1/2-e_u\sin^2\theta_W)(u_1+u_3)}{\cos\theta_W}V^*_{31}N_{i2}\right]\}$	$g^2\|K_{td}\|\left[\frac{(m_um_dd_1N^*_{i3}U_{32}+m_bm_td_3N_{i3}U^*_{32})}{2M_W^2\cos^2\beta}\right.$ $\left.+\frac{(m_um_dd_1N^*_{i4}V_{32}+m_bm_td_3N_{i4}V^*_{32})}{2M_W^2\sin^2\beta}\right]$	4.8×10^{-3}	$b\to ue\nu_i$
$-\sqrt{2}\|K_{td}\|\,\{e[V_{41}e_dN^*_{i1}(d_1+d_3)$ $-V^*_{41}e_uN_{i1}(u_1+u_3)]$ $-g\left[\frac{(1/2+e_d\sin^2\theta_W)}{\cos\theta_W}N^*_{i2}\right](d_1+d_3)$ $-\frac{(1/2-e_u\sin^2\theta_W)}{\cos\theta_W}N_{i2}](u_1+u_3)\Big]\}$	$g^2\|K_{td}\|\left[\frac{(m_um_dd_1N^*_{i3}U_{42}+m_bm_td_3N_{i3}U^*_{42})}{2M_W^2\cos^2\beta}\right.$ $\left.+\frac{(m_um_dd_1N^*_{i4}V_{42}+m_bm_td_3N_{i4}V^*_{42})}{2M_W^2\sin^2\beta}\right]$	5.3×10^{-3}	$b\to u\mu\nu_i$
$-\sqrt{2}\|K_{td}\|\,\{e\,[V_{51}e_dN^*_{i1}(d_1+d_3)$ $-V^*_{51}e_uN_{i1}(u_1+u_3)]$ $-g\left[\frac{(1/2+e_d\sin^2\theta_W)}{\cos\theta_W}N^*_{i2}(d_1+d_3)\right.$ $-\frac{(1/2-e_u\sin^2\theta_W)}{\cos\theta_W}N_{i2}](u_1+u_3)\Big]\}$	$g^2\|K_{td}\|\left[\frac{(m_um_dd_1N^*_{i3}U_{52}+m_bm_td_3N_{i3}U^*_{52})}{2M_W^2\cos^2\beta}\right.$ $\left.+\frac{(m_um_dd_1N^*_{i4}V_{52}+m_bm_td_3N_{i4}V^*_{52})}{2M_W^2\sin^2\beta}\right]$	3.1×10^{-3}	$b\to u\tau\nu_i$
$g^2(u_2+u_3)V^*_{31}V_{31}\|K_{cb}\|$	$\frac{g^2\|K_{cb}\|(m_s^2u_2+m_b^2u_3)}{2M_W^2\cos^2\beta}U^*_{32}U_{32}$	4.3×10^{-4}	$b\to se^+e^-$
$g^2(u_2+u_3)V^*_{41}V_{41}\|K_{cb}\|$	$\frac{g^2\|K_{cb}\|(m_s^2u_2+m_b^2u_3)}{2M_W^2\cos^2\beta}U^*_{42}U_{42}$	4.4×10^{-4}	$b\to s\mu^+\mu^-$
$g^2(u_2+u_3)(V^*_{31}V_{41}+V^*_{41}V_{31})\|K_{cb}\|$	$\frac{g^2\|K_{cb}\|(m_s^2u_2+m_b^2u_3)}{2M_W^2\cos^2\beta}$ $\times(U^*_{32}U_{42}+U^*_{42}U_{32})$	2.7×10^{-4}	$b\to se^\pm\mu^\mp$

One obtains similar constraints from semileptonic decays of the K meson, $K\to\pi l^+l^-$ and $K\to\pi\nu\bar{\nu}$. The experimental data on these decays is [13]:

$$
\begin{aligned}
BR(K\to\pi e^+e^-) &= (2.88\pm0.13)\times10^{-7},\\
BR(K\to\pi\mu^+\mu^-) &= (7.6\pm2.1)\times10^{-7},\\
BR(K\to\pi\mu^+e^-) &< 2.1\times10^{-10},\\
BR(K\to\pi e^+\mu^-) &= 7\times10^{-9}.
\end{aligned}
\qquad (32)
$$

With spontaneous R-parity violation, these decays can proceed at tree-level through either a $\tilde{u}$ or $\tilde{c}$ exchange, giving the constraints in Table 2.

The decay $K \to \pi\nu\bar{\nu}$ has an extremely small branching ratio [13]:

$$BR(K \to \pi\nu\bar{\nu}) = (1.5^{+3.4}_{-1.2}) \times 10^{-10} \tag{33}$$

These decay can proceed at tree-level through either the exchange of a $\tilde{d}$ or $\tilde{s}$, with the bounds in Table 2.

Finally, we look at the semileptonic decays of the D-meson, $D \to \bar{K}^{0*}e^+\nu_e$ and $D \to \bar{K}^{0*}\mu^+\nu_\mu$. The experimental inputs used are [13]:

$$\begin{aligned} BR(D \to \bar{K}^{0*}e^+\nu_e) &= (4.8 \pm 0.5) \times 10^{-2} \\ BR(D \to \bar{K}^{0*}e^+\nu_e) &= (4.4 \pm 0.6) \times 10^{-2} \end{aligned} \tag{34}$$

These decays can occur at tree-level though the exchange of either a $\tilde{s}$ or $\tilde{c}$ and we derive the constraints given in Table 2.

Table 2: Analytic bounds on mixing matrices of chargino-leptons $U_{ij},\ V_{ij}$ (with $j = 3, 4, 5$ corresponding to $e,\ \mu$ and τ) and neutralino-neutrinos N_{ij} (with $j = 5, 6, 7$ corresponding to ν_e, ν_μ and ν_τ) from K and D rare semileptonic decays.

gaugino type	higgsino type	bound	process
$g^2 V_{31}V^*_{31}\lvert K_{us}\rvert(u_1+u_2)$	$\frac{g^2\lvert K_{us}\rvert(m_u^2 u_1 + m_c^2 u_2)}{2M_W^2\cos^2\beta}U^*_{32}U_{32}$	1.4×10^{-4}	$K^+ \to \pi^+e^+e^-$
$g^2 V_{41}V^*_{41}\lvert K_{us}\rvert(u_1+u_2)$	$\frac{g^2\lvert K_{us}\rvert(m_u^2 u_1 + m_c^2 u_2)}{2M_W^2\cos^2\beta}U^*_{42}U_{42}$	1.4×10^{-4}	$K^+ \to \pi^+\mu^+\mu^-$
$g^2(V_{31}V^*_{41} + V^*_{31}V_{41})\lvert K_{us}\rvert(u_1+u_2)$	$\frac{g^2\lvert K_{us}\rvert(m_u^2 u_1 + m_c^2 u_2)}{2M_W^2\cos^2\beta}$ $\times(U^*_{32}U_{42} + U^*_{42}U_{32})$	1.4×10^{-4}	$K^+ \to \pi^+(e^+\mu^- +\mu^+e^-)$
$[ee_dN_{51} - \frac{g(1/2+e_d\sin^2\theta_W)}{\cos\theta_W}N_{52}]^2$ $\times 2\lvert K_{us}\rvert(d_1+d_2)$	$\frac{g^2\lvert K_{us}\rvert(m_d^2 d_1 + m_s^2 d_2)}{2M_W^2\cos^2\beta}N^*_{53}N_{53}$	1.6×10^{-5}	$K \to \pi\nu\bar{\nu}$
$-\sqrt{2}(V_{31}[ee_dN_{51} - \frac{g(1/2+e_d\sin^2\theta_W)}{\cos\theta_W}N_{52}]d_2$ $+V_{31}[ee_uN_{51} + \frac{g(1/2-e_d\sin^2\theta_W)}{\cos\theta_W}N_{52}]u_2)$	$\frac{g^2 m_c m_d}{2M_W^2\cos\beta\sin\beta}$ $\times(V^*_{32}N_{53}d_2 + U_{32}N^*_{54}u_2)$	0.13	$D \to \bar{K}^{0*}e^+\nu_e$
$-\sqrt{2}(V_{41}[ee_dN_{51} - \frac{g(1/2+e_d\sin^2\theta_W)}{\cos\theta_W}N_{62}]d_2$ $+V_{41}[ee_uN_{61} + \frac{g(1/2-e_d\sin^2\theta_W)}{\cos\theta_W}N_{62}]u_2)$	$\frac{g^2 m_c m_d}{2M_W^2\cos\beta\sin\beta}$ $\times(V^*_{42}N_{63}d_2 + U_{42}N^*_{64}u_2)$	0.09	$D \to \bar{K}^{0*}\mu^+\nu_\mu$

3.2 Rare leptonic decays

We now turn to the analysis of decays of the τ or μ. We denote the three-body leptonic decays of the μ or τ by $l_i \to l_j l_k l_m$, where i, j, k, m are generation indeces. The experimental bounds on these lepton flavor violating decays are [13]:

$$\begin{aligned}
BR(\mu^- \to e^- e^- e^+) &< 1 \times 10^{-12}, \\
BR(\tau^- \to e^+ \mu^- \mu^-) &< 2.9 \times 10^{-6}, \\
BR(\tau^- \to e^- e^- e^+) &< 1.5 \times 10^{-6}, \\
BR(\tau^- \to e^- \mu^- \mu^+) &< 1.8 \times 10^{-6}, \\
BR(\tau^- \to e^- e^- \mu^+) &< 1.5 \times 10^{-6}, \\
BR(\tau^- \to \mu^- \mu^- \mu^+) &< 1.9 \times 10^{-6}, \\
BR(\tau^- \to \mu^- e^- e^+) &< 1.7 \times 10^{-6}. \qquad (35)
\end{aligned}$$

These decays proceed by an exchange of a sneutrino $\tilde{\nu}_i$. We have tabulated the constraints, both for the gaugino- and higgsino-type couplings in Table 3.

We turn next to $\mu - e$ conversion. The conversion in nuclei is one of the most restricted leptonic phenomena. The upper limits extracted at PSI by the SINDRUM II experiments are [14, 15]:

$$R_{\mu e^-} < 6.1 \times 10^{-13} \quad \text{for } {}^{48}Ti \text{ target} \qquad (36)$$

$$R_{\mu e^-} < 4.6 \times 10^{-11} \quad \text{for } {}^{208}Pb \text{ target} \qquad (37)$$

This process is governed by the following effective Lagrangian:

$$\mathcal{L}_{eff} = \frac{1}{2}\bar{e}_L \gamma_\alpha \mu_L \left[A^d_{\mu T i} \bar{d}_R \gamma_\alpha d_R + A^u_{\mu T i} \bar{u}_R \gamma_\alpha u_R \right] + \frac{1}{2}\left[S^{d,1}_{\mu T i} \bar{e}_L \mu_R \bar{d}_R d_L + S^{d,2}_{\mu T i} \bar{e}_R \mu_L \bar{d}_L d_R \right] \qquad (38)$$

It can occur at tree-level through $\tilde{d}$ or $\tilde{\nu}$ quark exchange and it provides one of the most stringent bounds on mixing elements, as seen in Table 3.

The effective Lagrangian for muonium $M - \bar{M}$ conversion has a $(V - A) \times (V - A)$ structure as in the original papers [16]:

$$\mathcal{H} = G_{M\bar{M}} \bar{\psi}_\mu \gamma_\lambda (1 - \gamma_5) \psi_e \bar{\psi}_\mu \gamma_\lambda (1 - \gamma_5) \psi_e \qquad (39)$$

where the constant $G_{M\bar{M}}$ contains information on physics beyond the Standard Model. In our case, the process $\mu^+ e^- \to \mu^- e^+$, forbidden in the Standard Model, can proceed through tree-level graphs with either $\tilde{\nu}$ or $\tilde{d}$ exchanges (bounds obtained are in Table 3).

Next we investigate constraints coming from lepton family number violating decays of tau into a meson and a lepton, $\tau \to l + PS$ (or V), where $l = e$ or μ, $PS = \pi^0, \eta$, or K^0, and $V = \rho^0, \omega, K^*$, or ϕ. The amplitude obtained from the effective Lagrangians is:

$$\begin{aligned} \mathcal{M}(\tau \to l_k + V) &= \frac{1}{8} A_V f_V m_V \epsilon_\mu^* \bar{l}_k \gamma^\mu (1-\gamma_5)\tau \\ \mathcal{M}(\tau \to l_k + PS) &= \bar{l}_k (A_L^{PS} P_L + A_R^{PS} P_R)\tau \end{aligned} \tag{40}$$

The bounds on the gaugino and higgsino couplings come from the experimental data on the corresponding decays [13]:

$$\begin{aligned} BR(\tau^- \to e^-\pi^0) &< 3.7\times 10^{-6}, \\ BR(\tau^- \to \mu^-\pi^0) &< 4.0\times 10^{-6}, \\ BR(\tau^- \to e^-K^0) &< 1.3\times 10^{-3}, \\ BR(\tau^- \to \mu^-K^0) &< 1.0\times 10^{-3}, \\ BR(\tau^- \to e^-\eta) &< 8.2\times 10^{-6}, \\ BR(\tau^- \to \mu^-\eta) &< 9.6\times 10^{-6}, \\ BR(\tau^- \to e^-\rho^0) &< 2.0\times 10^{-6}, \\ BR(\tau^- \to \mu^-\rho^0) &< 6.3\times 10^{-6}, \\ BR(\tau^- \to e^-K^{0*}) &< 5.1\times 10^{-6}, \\ BR(\tau^- \to \mu^-K^{0*}) &< 7.4\times 10^{-6}. \end{aligned} \tag{41}$$

Both of these types of decays occur at tree-level through a $\tilde{u}$ or a $\tilde{d}$ exchange. As seen in Table 3, the constraints from $\tau \to l_k\phi$ are too weak to give any significant bounds on R-violating couplings.

Table 3: Analytic bounds on mixing matrices of chargino-leptons $U_{ij},\ V_{ij}$ (with $j=3,4,5$ corresponding to $e,\ \mu$ and τ) and neutralino-neutrinos N_{ij} (with $j=5,6,7$ corresponding to ν_e,ν_μ and ν_τ) from rare leptonic decays.

gaugino type	higgsino type	bound	process
$g^2V^*_{41}V_{31}n_1$	$\frac{g^2m_e^2}{2M_W^2\cos^2\beta}U^*_{42}V_{32}n_1$	6.6×10^{-7}	$\mu\to3e$
$g^2V^*_{51}V_{41}n_2$	$\frac{g^2m_\mu^2}{2M_W^2\cos^2\beta}U^*_{52}V_{42}n_2$	6.4×10^{-3}	$\tau\to3\mu$
$g^2V^*_{51}V_{31}n_1$	$\frac{g^2m_e^2}{2M_W^2\cos^2\beta}U^*_{52}V_{32}n_1$	5.6×10^{-3}	$\tau\to3e$
$g^2V^*_{51}V_{41}n_1$	$\frac{g^2m_e^2}{2M_W^2\cos^2\beta}U^*_{52}V_{42}n_1$	5.7×10^{-3}	$\tau\to2e\mu$
$g^2V^*_{51}V_{41}n_2$	$\frac{g^2m_\mu^2}{2M_W^2\cos^2\beta}V^*_{51}V_{42}n_2$	6.2×10^{-3}	$\tau\to2\mu e$
$g^2V^*_{41}V_{31}d_1$	$\frac{g^2m_u^2}{2M_W^2\sin^2\beta}V_{42}V^*_{32}n_1$	6.2×10^{-7}	$\mu-e$
$g^2(V^*_{41}V_{41}n_1+V^*_{31}V_{31}n_2)$	$\frac{g^2}{2M_W^2\cos^2\beta}(m_3^2U^*_{42}U_{42}n_1+m_\mu^2U^*_{32}U_{32}n_2)$	6.3×10^{-3}	$M-\bar{M}$
$g^2V^*_{51}V_{31}(d_1+u_1)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{32}m_u^2d_1+\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{32}m_d^2u_1$	3.5×10^{-3}	$\tau\to e\rho$
$g^2V^*_{51}V_{31}\|K_{us}\|(u_1+u_2)$	$\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{32}\|K_{us}\|(m_d^2u_1+m_s^2u_2)$	3.0×10^{-3}	$\tau\to eK^{0*}$
$g^2V^*_{51}V_{41}(d_1+u_1)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{42}m_u^2d_1+\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}m_d^2u_1$	4.2×10^{-3}	$\tau\to\mu\rho$
$g^2V^*_{51}V_{41}\|K_{us}\|(u_1+u_2)$	$\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}\|K_{us}\|(m_d^2u_1+m_s^2u_2)$	3.8×10^{-3}	$\tau\to\mu K^{0*}$
$g^2V^*_{51}V_{31}(d_1+u_1)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{32}m_u^2d_1-\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{32}m_d^2u_1$	6.6×10^{-2}	$\tau\to e\pi^0$
$g^2V^*_{51}V_{31}\|K_{us}\|(u_1+u_2)$	$\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{32}\|K_{us}\|(m_d^2u_1+m_s^2u_2)$	4.0×10^{-1}	$\tau\to eK^0$
$g^2V^*_{51}V_{41}(d_1+u_1)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{42}m_u^2d_1-\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}m_d^2u_1$	3.7×10^{-2}	$\tau\to\mu\pi^0$
$g^2V^*_{51}V_{31}(d_1+u_1-2u_2)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{42}m_u^2d_1$ $+\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}(m_d^2u_1-2m_s^2u_2)$	7.8×10^{-2}	$\tau\to e\eta$
$g^2V^*_{51}V_{41}\|K_{us}\|(u_1+u_2)$	$\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}\|K_{us}\|(m_d^2u_1+m_s^2u_2)$	3.6×10^{-1}	$\tau\to\mu K^0$
$g^2V^*_{51}V_{41}(d_1+u_1-2u_2)$	$\frac{g^2}{2M_W^2\sin^2\beta}V^*_{52}V_{42}m_u^2d_1$ $+\frac{g^2}{2M_W^2\cos^2\beta}V^*_{52}V_{42}(m_d^2u_1-2m_s^2u_2)$	8.2×10^{-2}	$\tau\to\mu\eta$

In the Table 4 below, we summarize our restrictions on R-violating mixing matrix elements from tree-level processes for a set of values of soft masses and $\tan\beta$.

Table 4: Numerical bounds on mixing matrices of chargino-leptons U_{ij}, V_{ij} and neutralino-neutrinos N_{ij} from rare decays for $m_{\tilde{f}} = 100$ GeV. Quark masses have been taken as $m_u = 5$ MeV, $m_d = 10$ MeV, $m_c = 1.5$ GeV, $m_s = 200$ MeV, and $m_b = 4.5$ GeV; and $\tan\beta = 2$.

Bound	Process	type
$\lvert V_{31}\rvert < 0.027,\ 0.113$	$K \to \pi e^+e^-$, $b \to s e^+e^-$	g
$\lvert U_{32}\rvert < 0.234,\ 0.304$	$K \to \pi e^+e^-$, $b \to s e^+e^-$	h
$\lvert V_{41}\rvert < 0.027,\ 0.113$	$K \to \pi\mu^+\mu^-$, $b \to s\mu^+\mu^-$	g
$\lvert U_{42}\rvert < 0.234,\ 0.308$	$K \to \pi\mu^+\mu^-$, $b \to s\mu^+\mu^-$	h
$Re(V_{31}^* V_{41}) < 0.014,\ 0.008$	$K \to \pi\mu^\pm e^\pm$, $b \to s\mu^\pm e^\pm$	g
$\lvert U_{32}U_{42}^*\rvert < .027,\ 0.029$	$K \to \pi\mu^\pm e^\pm$, $b \to s\mu^\pm e^\pm$	h
$Re(V_{51}^* V_{31}) < 0.686$	$B_d \to e^\pm\tau^\pm$	g
$Re(V_{51}^* V_{41}) < 0.876$	$B_d \to \mu^\pm\tau^\pm$	g
$\lvert V_{41}^* V_{31}\rvert < 1.56 \times 10^{-6},\ 1.67 \times 10^{-7}$	$\mu \to 3e,\ \mu - e$ conversion	g
$\lvert V_{51}^* V_{41}\rvert < 0.015,\ 0.0135,\ 0.004,\ 0.02$	$\tau \to 3\mu,\ \tau \to 2\mu e,\ \tau \to \mu\rho,\ \tau \to \mu K^{0*}$	g
$\lvert V_{51}^* V_{31}\rvert < 0.013,\ 0.015,\ 0.004$	$\tau \to 3e,\ \tau \to \mu 2e,\ \tau \to e\rho$	g
$\lvert 0.31N_{51} + 0.94N_{52}\rvert < 0.07$	$K \to \pi\nu\nu$	g
$\lvert N_{53}\rvert < 0.56$	$K \to \pi\nu\nu$	h
$\lvert V_{31}^*(0.93N_{51} + 1.71N_{52})\rvert < .786,\ .045,\ .425$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu},\ D \to \bar{K}^{0*}\bar{e}\nu_e$	g
$\lvert N_{53}U_{32}^* + N_{54}^* V_{32}\rvert < 0.116,\ .0066$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu}$	h
$\lvert V_{31}^*(0.93N_{61} + 1.71N_{62})\rvert < .786,\ .045,$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu}$	g
$\lvert N_{63}U_{32}^* + N_{64}^* V_{32}\rvert < 0.116,\ .0066$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu}$	h
$\lvert V_{31}^*(0.93N_{71} + 1.71N_{72})\rvert < .786,\ .045$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu}$	g
$\lvert N_{73}U_{32}^* + N_{74}^* V_{32}\rvert < 0.116,\ .0066$	$b \to ue\bar{\nu},\ b \to ce\bar{\nu}$	h
$\lvert V_{41}^*(0.93N_{51} + 1.71N_{52})\rvert < .868,\ .045$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu}$	g
$\lvert N_{53}U_{42}^* + N_{54}^* V_{42}\rvert < 0.358\ ,.0066$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu}$	h
$\lvert V_{41}^*(0.93N_{61} + 1.71N_{62})\rvert < .868,\ .045,\ .295$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu},\ D \to \bar{K}^{0*}\bar{\mu}\nu_\mu$	g
$\lvert N_{63}U_{42}^* + N_{64}^* V_{42}\rvert < 0.358,\ .0066$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu}$	h
$\lvert V_{41}^*(0.93N_{71} + 1.71N_{72})\rvert < .868,\ .045$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu}$	g
$\lvert N_{73}U_{42}^* + N_{74}^* V_{42}\rvert < 0.358\ ,.0066$	$b \to u\mu\bar{\nu},\ b \to c\mu\bar{\nu}$	h
$\lvert V_{51}^*(0.93N_{51} + 1.71N_{52})\rvert < .505,\ .045$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	g
$\lvert N_{53}U_{52}^* + N_{54}^* V_{52}\rvert < 0.274\ ,.0066$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	h
$\lvert V_{51}^*(0.93N_{61} + 1.71N_{62})\rvert < .505,\ .045,$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	g
$\lvert N_{63}U_{52}^* + N_{64}^* V_{52}\rvert < 0.274,\ .0066$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	h
$\lvert V_{51}^*(0.93N_{71} + 1.71N_{72})\rvert < .505,\ .045$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	g
$\lvert N_{73}U_{52}^* + N_{74}^* V_{52}\rvert < 0.274,.0066$	$b \to u\tau\bar{\nu},\ b \to c\tau\bar{\nu}$	h

(By "h" and "g" we mean higgsino or gaugino coupling.)

4 One Loop Processes

In addition to processes that can occur at tree-level, there are others which can only occur at one loop-level, but are highly suppressed; or processes like $\mu - e$ conversion, which may set more stringent limits at one-loop level than at tree-level. Usually these processes invove chirality flip on an internal or external leg. These are: the anomalous magnetic moment of the muon $(g-2)_\mu$, lepton flavor violating processes $\mu - e$ conversion and $\mu \to e\gamma$ and the lepton and quark electric dipole moments.

First we investigate the effect of spontaneous R-parity breaking on the decay $\mu \to e\gamma$. The amplitude of the $\mu \to e\gamma$ transition can be written in the form of the usual dipole-type interaction:

$$\mathcal{M}_{\mu\to e\gamma} = \frac{1}{2}\bar{\psi}_e(d_L P_L + d_R P_R)\sigma^{\mu\nu}F_{\mu\nu}\psi_\mu \tag{42}$$

It leads to the branching ratio:

$$BR(\mu \to e\gamma) = \frac{1}{16\pi}\tau_\mu(|d|_L^2 + |d|_R^2)m_\mu^3 \tag{43}$$

Comparing it with the standard decay width, $\Gamma_{\mu\to e\nu\bar{\nu}} = \frac{1}{192\pi^3}G_F^2 m_\mu^5$ and using the experimental constraint on the branching ratio $B.R.(\mu \to e\gamma) < 1.2 \times 10^{-11}$ [13], one obtains the following limit on the dipole amplitude:

$$|d| = \sqrt{(|d_L|^2 + |d_R|^2)/2} < 1.73 \cdot 10^{-26}\ e \cdot cm \tag{44}$$

A non-vanishing dipole interaction results in a fermion chirality flip. There are two possibilities for this to occur. One is that the chirality flip occurs on the external muon line, resulting in a proportionality of the decay amplitude to the muon mass. The other is that the chirality flip occurs on the internal line, resulting in proportionality of the same amplitude to the mass of the fermion in the loop. This latter process requires the mixing of the left and right squarks or sleptons, and the resulting amplitude is proportional to the mixing angle. In R-parity conserving SUSY, the latter process dominates due to the large fermion mass in the loop, and also due to the loop function which is larger (by an order of magnitude or more) than the corresponding one for the process with external chirality flip. The same is true with spontaneous R-parity breaking and the bound obtained is:

$$\frac{1}{16\pi^2}\frac{g^2|V_{j1}U_{j2}^*|}{2M_W\cos\beta}m_\mu\frac{f_1(x)}{m_{\tilde{f}_1}^2} < 1.47 \times 10^{-3} \tag{45}$$

where the loop integral is:

$$f_1(x) = \frac{1}{2(1-x)^2}\left[3 - x + \frac{2\ln x}{1-x}\right] \tag{46}$$

and $x = \frac{m_\tau^2}{m_{\tilde{f}}^2}$.

$\mu - e$ conversion in nuclei is perhaps the most interesting lepton-flavor violating process experimentally. From a theoretical point of view, it is the most difficult to disentangle, because of the interdependence between particle and nuclear physics elements, in particular the difficulty in evaluating nuclear matrix elements. The process is very interesting at one-loop level from two points of view. First, it has quite a different structure from $\mu \to e\gamma$ (as opposed to $\mu^+ \to e^+e^+e^-$). Therefore it provides complimentary information on muon decay from the first two decays: it can occur even when $\mu \to e\gamma$ is forbidden, and it could be a better indicator of a rich gauge structure, such as extra Z or W bosons. Second, it has been shown that for a class of models $\mu - e$ conversion is enhanced with respect to $\mu \to e\gamma$ by large $\ln(m_\mu^2/\Lambda^2)$, where Λ is the scale responsible for the new physics [17]. With the expected improvement in experimental data, this test is likely to become the most stringent in R-parity violation. Based on the above transition elements, the branching ratio for the coherent $\mu^- - e^-$ conversion is given by [18]:

$$\begin{aligned} R_{ph}(\mu^- N \to e^- N) &= \frac{p_e E_e Z \alpha^5 Z_{eff}^4 F_p^2}{m_\mu \Gamma_{capt}} \{|f_{E0}(-m_\mu^2) + f_{M1}(-m_\mu^2) + f_{M0}(-m_\mu^2) + f_{E1}(-m_\mu^2)|^2 \\ &+ |f_{E0}(-m_\mu^2) + f_{M1}(-m_\mu^2) - f_{M0}(-m_\mu^2) - f_{E1}(-m_\mu^2)|^2\} \end{aligned} \tag{47}$$

where Γ_{capt} is the total muon capture rate, Z_{eff} is an effective atomic charge obtained by averaging the muon wave function over the nuclear density, and F_p is the nuclear matrix element. The functions f_{E0}, f_{E1}, f_{M0} and f_{M1} depend on loop functions and on the R-parity violating couplings [17]. The bounds obtained, listed in Table 4, restrict the same combination of parameters as the bounds obtained from $\mu \to e\gamma$ but slightly stronger. As for the radiative decays of the τ lepton, $\tau \to \mu\gamma$ and $\tau \to e\gamma$, they do not add anything new to the the bounds found so far. These radiative decays constrain the same combination of mixing matrix elements, but the bounds are much weaker, owing to weaker experimental limits of the radiative decays of the τ versus the μ.

Next we evaluate the contributions coming from the electric dipole moments. The electric dipole moment of an elementary fermion is defined through its electromagnetic form factor

$F_3(q^2)$ found from the (current) matrix element:

$$\langle f(p')|J_\mu(0)|f(p)\rangle = \bar{u}(p')\Gamma_\mu(q)u(p), \tag{48}$$

where $q = p' - p$ and

$$\Gamma_\mu(q) = F_1(q^2)\gamma_\mu + F_2(q^2)i\sigma_{\mu\nu}q^\nu/2m + F_A(q^2)(\gamma_\mu\gamma_5 q^2 - 2m\gamma_5 q_\mu) + F_3(q^2)\sigma_{\mu\nu}\gamma_5 q^\nu/2m, \tag{49}$$

with m the mass of the fermion. The EDM of the fermion field f is then given by

$$d_f = -F_3(0)/2m, \tag{50}$$

corresponding to the effective dipole interaction

$$\mathcal{L}_I = -\frac{i}{2}d_f\bar{f}\sigma_{\mu\nu}\gamma_5 f F^{\mu\nu} \tag{51}$$

The effective Lagrangian is induced at one-loop level if the theory contains a CP-violating coupling at tree-level. We can parametrize the interaction of a fermion Ψ_f with other fermions Ψ_i-s and scalars Φ_k-s, with respective charges Q_f, Q_i and Q_k, in general as:

$$-\mathcal{L}_{int} = \sum_{ik}\bar{\Psi}_f\left(A_{ik}\frac{1-\gamma_5}{2} + B_{ik}\frac{1+\gamma_5}{2}\right)\Psi_i\Phi_k + H.C. \tag{52}$$

If there is CP-violation, then $Im(A_{ik}B^*_{ik}) \neq 0$, and the one-loop fermion EDM is given by:

$$d^E_f = \sum_{ik}\frac{m_i}{(4\pi)^2 m_k^2}Im(A_{ik}B^*_{ik})\left(Q_i f_1(\frac{m_i^2}{m_k^2}) + Q_k f_2(\frac{m_i^2}{m_k^2})\right) \tag{53}$$

with:

$$f_2(x) = \frac{1}{2(1-x)^2}\left(1 + x + \frac{2x\ln x}{1-x}\right), \tag{54}$$

assuming charge conservation at the vertices $Q_k = Q_f - Q_i$. Since a non-vanishing d_f in the SM results in fermion chirality flip, it requires both CP violation and $SU(2)_L$ symmetry breaking. Experimentally, the EDMs of the electron and the neutron are some of the most restrictive parameters in the Particle Data Booklet, the present experimental upper limits being $d_e \leq 4.3\cdot 10^{-27}ecm$ and $d_n \leq 6.3\cdot 10^{-26}ecm$ [13].

The spontaneous R-violating contribution to the dipole moment of an electron is:

$$d^E_e = \frac{\alpha_{EM}}{4\pi\sin^2\theta_W}Im(V_{i1}U_{i2})\frac{m_e m_{e_i}}{\sqrt{2}M_W\cos\beta}\frac{f_1(x_e)}{m^2_{\tilde{f}}} < 4.3\times 10^{-27}ecm \tag{55}$$

with $x_e = \frac{m_{e_i}^2}{m_{\tilde{f}}^2}$ and the d-quark contribution is:

$$
\begin{aligned}
d_d^E \; &= N_c \frac{\alpha_{EM}}{4\pi \sin^2\theta_W} \left\{ Im(V_{i1}U_{i2}) \frac{m_d m_{e_i}}{\sqrt{2} M_W \cos\beta} \frac{f_1(x_i)}{m_{\tilde{f}}^2} + \sqrt{2} \left[e e_d Im(N_{71}N_{73}) \right.\right. \\
- \; & \left.\left. \frac{g}{\cos\theta_W} (-\frac{1}{2} - e_d \sin^2\theta_W) Im(N_{72}N_{73}) \right] \frac{F(x_\nu)}{m_{\tilde{d}}^2} \right\} < 4.725 \times 10^{-26} ecm
\end{aligned} \tag{56}
$$

with $F = f_1 + 2 f_2$, $x_\nu = \frac{m_\nu^2}{m_{\tilde{f}}^2}$. To evaluate the EDM of the neutron we use:

$$
d_n = \frac{4}{3} d_d - \frac{1}{3} d_u \tag{57}
$$

We include for completeness the constraint arising from the new measurement of the anomalous magnetic moment of the muon. The new measurement for the muon anomalous magnetic moment a_μ corresponds to a deviation from the Standard Model prediction:

$$
a_\mu^{exp} - a_\mu^{SM} = (4.26 \pm 1.65) \times 10^{-9} \tag{58}
$$

If the deviation can be attributed to new physics effects, then at 90% C.L. δa_μ^{NP} must lie in the range:

$$
2.15 \times 10^{-9} \le \delta a_\mu^{NP} \le 6.37 \times 10^{-9} \tag{59}
$$

The anomalous magnetic moment of the muon arises from terms of the form:

$$
\frac{ie}{2m_\mu} F(q^2) \bar{\psi} \sigma_{\alpha\beta} q^\beta \psi \tag{60}
$$

with $a_\mu = F(0)$. The contributions to a_μ are proportional to the mass of the muon squared:

$$
\delta a_\mu = \frac{m_\mu^2}{2} (A_L^{22} + A_R^{22}) \tag{61}
$$

From spontaneous R-parity violation, we obtain the bound:

$$
\frac{1}{4\pi^2} \frac{g^2 m_\mu^2}{2 M_W \cos\beta} N_{i3} [\sin\theta_W N_{i1} - \frac{g}{\cos\theta_W} (\frac{1}{2} - \sin^2\theta_W) N_{i2}] \frac{f_2(x_\nu)}{m_{\tilde{f}}} < 4.2 \times 10^{-9} \tag{62}
$$

where $i = 5, 6$ or 7. We take $m_{\nu_\tau} = 1$ eV.

In Table 4 below we summarize all one-loop bounds we obtained. In the case in which more than a term is present in a constraint, and we have insufficient information to bound the terms separately, we obtain the bounds by assuming that only one term is non-zero. All of these bounds include products of couplings from vertices including higgsino or gaugino, and are thus new bounds, not present in models with explicit trilinear R-parity violation.

Table 4: Numerical bounds on mixing matrices of chargino-leptons U_{ij}, V_{ij} and neutralino-neutrinos N_{ij} from one-loop processes for $m_{\tilde{f}} = 100$ GeV and $\tan\beta = 2$.

Bound	Process	type
$Re(V^*_{31}U^*_{32}) < 2 \times 10^{-5}$	$\mu \to e\gamma$	g, h
$Re(V^*_{41}U^*_{42}) < 3.8 \times 10^{-5}$	$\mu \to e\gamma$	g, h
$Re(V^*_{51}U^*_{52}) < 6.8 \times 10^{-5}$	$\mu \to e\gamma$	g, h
$Re(N_{i1}N_{i3}) < 1.4 \times 10^{-2}$	$\mu \to e\gamma$	g, h
$Re(N_{i2}N_{i3}) < 8.5 \times 10^{-2}$	$\mu \to e\gamma$	g, h
$Re(V^*_{31}U^*_{32}) < 3 \times 10^{-3}$	$(g-2)_\mu$	g, h
$Re(V^*_{41}U^*_{42}) < 6.1 \times 10^{-3}$	$(g-2)_\mu$	g, h
$Re(V^*_{51}U^*_{52}) < 1.1 \times 10^{-2}$	$(g-2)_\mu$	g, h
$Re(V^*_{31}U^*_{32}) < 5.6 \times 10^{-6}$	$\mu - e$ conversion	g, h
$Re(V^*_{41}U^*_{42}) < 1.1 \times 10^{-5}$	$\mu - e$ conversion	g, h
$Re(V^*_{51}U^*_{52}) < 2 \times 10^{-5}$	$\mu - e$ conversion	g, h
$Re(N_{i1}N_{i3}) < 9.3 \times 10^{-3}$	$\mu - e$ conversion	g, h
$Re(N_{i2}N_{i3}) < 5.7 \times 10^{-3}$	$\mu - e$ conversion	g, h
$Im(V_{41}U_{42}) < 6.2 \times 10^{-2}$	EDM_e	g, h
$Im(V_{51}U_{52}) < 6.9 \times 10^{-3}$	EDM_e	g, h
$Im(V_{41}U_{42}) < 2.9 \times 10^{-3}$	EDM_n	g, h
$Im(V_{51}U_{52}) < 1.7 \times 10^{-4}$	EDM_n	g, h
$Im(N_{i1}N_{i3}) < 2.8 \times 10^{-2}$	EDM_n	g, h
$Im(N_{i2}N_{i3}) < 9.3 \times 10^{-3}$	EDM_n	g, h

where $i = 5, 6, 7$.

5 Conclusion

Conservation of R-parity, introduced to distinguish particles from their supersymmetric partners, is not imposed by any symmetry of the model. Explicit R-parity violation, allowed in MSSM, may not be allowed by higher gauge structures. However, the R-parity may be broken spontaneously through the Higgs mechanism. This type of breaking is achieved through vevs for the sneutrino fields. It has the attractive feature that it only breaks lepton number, thus avoiding fast proton decay. It allows for a dynamical mechanism to break R, much like electroweak symmetry breaking.

Spontaneous R parity breaking generates bilinear terms in the Lagrangian with both gaugino- and higgsino-type couplings. In this work, we assumed a general pattern of neutrino-neutralino and lepton-chargino mixing. Although the particle content of a given supersymmetric model will have to be enlarged to allow for spontaneous R-parity breaking, we deal

with a truncated version and assume an effective MSSM particle content. We then set general constraints on mixing matrix elements, valid for any supersymmetric model with spontaneous R parity breaking. For tree-level processes, we obtain some mass-dependent bounds and also some mass-independent bounds which come from gaugino-type couplings, most of which are new. Restricting processes which require chirality flip (at one-loop level), we obtain strong bounds on products of gaugino and higgsino couplings all of which are new. These results are complementary to those previously found [19].

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